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A Self-Consistent Model For The Long-Term Gamma-Ray Spectral Variability of Cygnus X-1

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ABSTRACT

The long-term transitions of the black hole candidate Cygnus X-1 (between the states γ_1 , γ_2 , and γ_3) include the occasional appearance of a strong \sim MeV bump (γ_1), whose strength appears to be anti-correlated with the continuum flux (\lesssim 400 keV) due to the Compton upscattering of cold disk photons by the inner, hot corona. We develop a self-consistent disk picture that accounts naturally for these transitions and their corresponding spectral variations. We argue that the bump is due to the self-Comptonization of bremsstrahlung photons emitted predominantly near the plane of the corona itself. Our results suggest that a decrease by a factor of ≈ 2 in the viscosity parameter α is responsible for quenching this bump and driving the system to the γ_2 state, whereas a transition from γ_2 to γ_3 appears to be induced by an increase of about 25% in the accretion rate \dot{M} . In view of the fact that most of the transitions observed in this source seem to be of the $\gamma_2 - \gamma_3$ variety, we conclude that much of the long term gamma-ray spectral variability in Cygnus X-1 is due to these small fluctuations in \dot{M} . The unusual appearance of the γ_1 state apparently reflects a change in the dissipative processes within the disk.

Subject headings: accretion disks—black hole physics—radiation mechanisms: bremsstrahlung—radiation mechanisms: inverse Compton—relativity—stars: individual: Cygnus X-1

1. Introduction

The black hole candidate Cygnus X-1 has shown variability of its X-ray emission on all time scales ranging from milliseconds to years. Early observations of this strong hard X-ray source (e.g., Tananbaum et al. 1972; Holt et al. 1976; Ogawara et al. 1982) revealed long-term transitions between two principal states: the so-called “low-state” with a relatively small ratio of soft to hard X-ray flux, and the “high-state” corresponding to a softer spectrum. Generally speaking, the hard X-ray/ γ -ray emission has been successfully interpreted in terms of the Compton upscattering of soft disk photons by an optically thin, high-temperature plasma ($\sim 10^9$ K) situated either at the inner edge of the disk (Thorne & Price 1975; Shapiro, Lightman & Eardley 1976), or within a coronal layer overlying the geometrically thin but optically thick cooler disk (Liang & Price 1977; Sunyaev & Titarchuk 1985). The first scenario could arise as a result of the secular instability present in the cool disk’s inner region (Lightman & Eardley 1974), which would swell the optically thick, radiation-pressure dominated gas to a hot, gas-pressure dominated, optically thin cloud.

More recent observations with the JPL High Resolution Gamma-Ray Spectroscopy experiment onboard *HEAO 3* (Ling et al. 1987) have suggested a more complicated behavior than that assumed in a simplified two-state model. In particular, a new γ -ray emitting state of Cygnus X-1 has been identified, in which the luminosity within an unusually strong γ -ray “bump” between ~ 400 keV and ~ 2 MeV exceeds that of the $50 - 400$ keV hard X-rays and is comparable to the overall emission below ~ 400 keV. This state, labeled γ_1 , together with the other low-state (dubbed γ_2) and the high-state (γ_3), seem to form a sequence of transitions triggered, perhaps, by variations in the accretion rate, as proposed earlier for the two-state system (e.g., Kazanas 1986).

Liang and Dermer (1988) have interpreted this γ -ray bump as due to the emission from a hot quasi-spherical pair-dominated cloud (distinct from the corona producing the continuum), in which the pair-balance condition determines the compactness parameter and the Thomson depth as a function of the equilibrium temperature. They find that

a reasonable fit to the “bump” may be obtained for a lepton temperature ≈ 400 keV and a pair cloud radius of about $9 r_g$, where $r_g \equiv 2GM/c^2$ is the Schwarzschild radius for a black hole mass M . Our purpose in this *Letter* is to build on the strengths of these earlier models to arrive at a self-consistent accretion disk picture that accounts naturally for these transitions and their corresponding spectral variations. In particular, we relax the assumption that the inner hot plasma is uniform (Shapiro, Lightman & Eardley 1976; Liang & Dermer 1988) and consider a more realistic stratified structure consistent with gravitational settling. As we shall see, the secular instability still results in a two-temperature corona in the inner region of the disk, but the electron density in the disk plane may now be sufficiently high to produce a copious supply of bremsstrahlung photons. We shall argue that the self-Comptonization of this spectral component is the origin of the bump seen in the γ_1 state. Not surprisingly, the physical conditions in this inner hot corona are not unlike those inferred for the pair plasma by Liang and Dermer (1988), but here the lepton density is due mostly to accretion, not pair production.

2. A Self-Consistent Two-Temperature Disk

2.1. The Structure Equations

Our starting point is to assume a standard α -disk model (Shakura & Sunyaev 1973), which however is subject to the secular instability in the inner region that drives the disk from a cool state (with electron temperature $T_e \sim 10^6$ K) to the hot, two-fluid state ($T_e \sim 10^9$ K and ion temperature $T_i \gg T_e$; Shapiro, Lightman & Eardley 1976, hereafter SLE). Unlike the earlier work, we do not employ vertically-averaged equations for the coronal structure, but rather allow for density stratification as discussed above. For simplicity, we assume constant temperatures in the vertical (i.e., z) direction, though not in the (cylindrical) radial direction (R). As such, SLE’s equations (3) and (4) are replaced by

$$\frac{dP_g}{dz} = -\frac{GM\rho z}{r^3} \quad (1)$$

and

$$\int_{-\infty}^{\infty} \alpha P_g(z) (2\pi R) R dz = (GMR)^{1/2} \dot{M} \mathcal{J} , \quad (2)$$

respectively, where $r^2 = R^2 + z^2$, P_g is the gas pressure, ρ is the density and $\mathcal{J} \equiv 1 - (3r_g/r)^{1/2}$. We ignore the contribution to the lepton number density from pair production and check *a posteriori* that this is indeed a valid approximation. In addition, the electron-ion energy exchange flux (SLE equation 10) generalizes to

$$\int dF = \int \frac{3}{2} \nu_E \rho k(T_i - T_e)/m_p dz , \quad (3)$$

where ν_E is the electron-ion coupling rate.

It is not difficult to see that with these modifications, the two-temperature structure of the hot inner corona is now specified by

$$\rho = \rho_0 \exp(-z^2/z_0^2) , \quad (4)$$

where

$$z_0 \approx (1.7 \times 10^5 \text{ cm}) \left[\frac{\mathcal{J}}{\alpha} \left(\frac{M}{3 M_\odot} \right) \left(\frac{\dot{M}}{10^{17} \text{ g s}^{-1}} \right) \right]^{1/3} \left(\frac{\rho_0}{5 \times 10^{-5} \text{ g cm}^{-3}} \right)^{-1/3} \left(\frac{2r}{r_g} \right)^{1/2} , \quad (5)$$

and the temperatures are given by

$$T_e + T_i \approx \frac{(8.3 \times 10^{11} \text{ K})}{(2r/r_g)^2} \left[\frac{\mathcal{J}}{\alpha} \left(\frac{\dot{M}}{10^{17} \text{ g s}^{-1}} \right) \right]^{2/3} \left(\frac{M}{3 M_\odot} \right)^{-4/3} \left(\frac{\rho_0}{5 \times 10^{-5} \text{ g cm}^{-3}} \right)^{-2/3} \quad (6)$$

$$\frac{T_e}{(T_i - T_e)^{2/3}} \approx \left[\frac{8103}{\alpha^2 \mathcal{J}^4} \left(\frac{M}{3 M_\odot} \right)^{14} \left(\frac{\dot{M}}{10^{17} \text{ g s}^{-1}} \right)^{-4} \left(\frac{\rho_0}{5 \times 10^{-5} \text{ g cm}^{-3}} \right)^{10} \left(\frac{2r}{r_g} \right)^{21} \right]^{1/9} . \quad (7)$$

This prescription, however, is inappropriate whenever the physical conditions are such that the protons attain sufficient energy to escape from the system. Since the virial temperature

$$T_{vir} \equiv \frac{1}{6} m_p c^2 \left(\frac{r_g}{r} \right) , \quad (8)$$

where m_p is the proton mass, decreases with increasing radius r , it is anticipated that the ion temperature T_i may exceed T_{vir} at large radii. In that case, a coronal mass (and energy) outflow will prevent T_i from greatly surpassing T_{vir} . For the purpose of this calculation, we shall assume that the fraction of dissipated power lost in this wind is at most a small fraction of the radiated luminosity. (This restriction will be relaxed in future radiative-hydrodynamical simulations of the accretion geometry.) To simplify matters further, we shall also assume that in this region (where the upper layers of the corona merge into a “wind”), the ion temperature is in fact equal to the virial temperature (though of course this can only be approximately correct). Instead of the above set of equations, the relevant expressions that determine the coronal structure are then Equations (1), (3), and the condition $T_i = T_{vir}$, which yield

$$\rho = \rho_0(1 + z^2/R^2)^{-3/2}, \quad (9)$$

and

$$T_e = 5.2 \times 10^9 \text{ K} \left(\frac{R}{r_g} \right)^2 \left(\frac{\rho_0}{5 \times 10^{-5} \text{ g cm}^{-3}} \right)^{4/3} \left(\frac{\dot{M}}{10^{17} \text{ g s}^{-1}} \right)^{-2/3} \left(\frac{M}{3 M_\odot} \right)^2. \quad (10)$$

In writing these equations, we have assumed that although the outer coronal layers may merge into a transonic flow (at say z_{sonic}), most of the structure of interest to us here lies at $z \ll z_{sonic}$, for which Equation (1) is still an adequate representation of the vertical density profile near the equatorial plane.

2.2. Method Of Solution

Following SLE, we have taken the outer boundary of the two-temperature inner region to lie at a radius R_0 given by the condition $P_r(R_0) = 3P_g(R_0)$, where P_r is the radiation pressure in the α -disk. We note, however, that our results are insensitive to the actual location of this boundary, as long as $P_r(R_0) \sim O[P_g(R_0)]$. For example, when R_0 is instead fixed by the condition $P_r(R_0) = 3/2 P_g(R_0)$, our spectra differ from those presented below by at most 10%, well within the observational uncertainties. In this model, the cooling results from the inverse Comptonization of both the cold disk photons penetrating into the

corona and the bremsstrahlung radiation produced in this inner hot region. To handle this dichotomy, we divide the corona into concentric rings and proceed as follows. We assume a fractional cooling by each spectral component in the first outer zone and then iterate over the mid-plane density until the local bremsstrahlung-self-Compton emissivity (within this ring) is balanced by the assumed fraction of the dissipation rate at this location. Within any given ring, the contribution to the overall flux from the Comptonized (outer-disk) photons is proportional to the total number of scattering particles and the electron temperature within that ring, and the attenuated soft photon number density, as reflected by the dependence of the Kompane'ets equation on these parameters. Thus, in subsequent rings, we iterate on the local mid-plane density until the calculated fraction agrees with the value extrapolated on the basis of this proportionality from the adjacent outer ring. From this we determine the overall coronal luminosity, including both components, and then iterate on the assumed fraction in the outer ring until this power matches the total dissipation rate inside the corona. We check ring by ring whether the first set of structure equations (see §2.1. above) yields an ion temperature $T_i > T_{vir}$. If mass loss is indicated at a particular radius, we then use the second set of equations to determine the local structure.

The cooling rate (and concomitantly the spectrum) is obtained by solving the relativistic steady-state Fokker-Planck equation, whose diffusion coefficient is valid for *arbitrary* photon and electron energies (Prasad et al. 1988). At the temperatures of interest in our problem, the use of the non-relativistic Kompane'ets equation (Kompane'ets 1956) is not valid, which in the past has led to the development of a (time-consuming) Monte-Carlo scheme in order to treat the inverse Compton scattering correctly (e.g., Liang & Dermer 1988). With the recent development of the exact analytical formula for this (relativistically-correct) diffusion coefficient of the Compton Fokker-Planck equation, we are now able to circumvent this difficulty by simply solving the relativistic Kompane'ets equation without any approximation. In this regard, our approach differs from that of SLE in at least 2 ways. First, our source term includes both the cool disk photons penetrating the $y \lesssim 1$ region of

the corona and the bremsstrahlung photons produced within the corona itself (predominantly near the plane). Here, $y \equiv (4kT_e/m_e c^2)\text{Max}(\tau, \tau^2)$ is the dimensionless parameter that characterizes Comptonization, and τ is the scattering optical depth. Secondly, since our equation is valid at all energies, our results are not subject to the restrictions (e.g., in T_e) and errors inherent in Cooper’s modification to the Kompane’ets equation (Cooper 1971). Of course, the main result of our calculations is the spectrum, and this is determined along the lines developed by SLE, in which the energy-dependent rate of photon loss appearing in the Fokker-Planck equation also acts as the source term for the energy-dependent photon flux.

3. The Gamma-Ray Spectral Variability

Since the structure of the hot inner corona is here calculated self-consistently, we have at most only 3 physical quantities that need to be fixed. These are the black hole mass M , the accretion rate \dot{M} , and the viscosity parameter α . A very conservative lower mass limit for Cygnus X-1 is $3.4 M_\odot$ (e.g., Paczyński 1974), whereas the most likely value is thought to lie in the range $9 - 15 M_\odot$ (e.g., Avni & Bahcall 1975). We have chosen to work with a value $M = 10 M_\odot$, though again our results are insensitive to the actual mass as long as $5 M_\odot \lesssim M \lesssim 15 M_\odot$. The basis for this “robustness” is the fact that distances scale as the Schwarzschild radius $r_g = 2GM/c^2$, so that the overall energy budget ($\sim G\dot{M}M/r$) is independent of M (to first order) for a given \dot{M} . Nonetheless, the coronal structure does change with M , but we have found only $\lesssim 10\%$ differences in the spectra over this mass range and have therefore not considered M to be a principal parameter characterizing our solutions. In addition, although we expect (slight) variations in \dot{M} between the various states of Cygnus X-1, the gross X-ray/ γ -ray energetics for an assumed source distance of 2.5 kpc restrict the accretion rate to a value $\approx 10^{18} \text{ g s}^{-1}$.

Our solutions for the γ_1 , γ_2 , and γ_3 states are compared to the data in Figures 1, 2, and 3. The data used by SLE overlap most significantly with the γ_1 spectrum and have therefore been included in Figure 1. We note also that we have excluded the 1.5 MeV point in the HEAO 3 data since it may be contaminated by a background line (Liang and Dermer

1988). In each case, the dashed curve shows the contribution from the bremsstrahlung/self-Comptonized (BSC) photons originating from within the corona itself, whereas the solid curve represents the overall spectrum comprising the BSC radiation, the cold disk photons that are Compton upscattered in the $y \lesssim 1$ region of the corona, and the thermal emission from the disk itself (which dominates at $E \lesssim 3$ keV). What emerges from these results is the interesting correlation between the variation in α and the gradual shift of emitted power from the γ -ray continuum into the broad \sim MeV bump (see Figure 4). In particular, a factor 2 decrease in α is sufficient to drive the system from the γ_1 to the γ_2 states. Thus, the emergence of the bump in the γ_2 and especially the γ_1 states must be accompanied by a corresponding reduction in the continuum flux below about 400 keV, as confirmed by the data. On the other hand, an increase in \dot{M} raises the flux in both the continuum and the bump. It appears that an increase of about 25% in \dot{M} is responsible for driving the system from the γ_2 to the γ_3 states (see Figure 5).

Figures 6, 7, 8, and 9 demonstrate how the internal structure of the corona changes in response to a variation in either α (characterizing the transition from the γ_1 to the γ_2 states) or the accretion rate (which induces the transition from γ_2 to γ_3). In Figure 6, the most significant aspect of the electron temperature profile is that the corona becomes optically thick toward small radii due to gravitational stratification (see the profiles of mid-plane density shown in Figure 7), and T_e increases rapidly in this region to enhance the bremsstrahlung-self-Compton cooling process. The profiles of ion temperature are shown in Figure 8. In all three states, T_i increases with radius until it reaches the virial temperature T_{vir} , though the radius at which this occurs is significantly larger in the γ_1 state than the other two. The impact of this change in the ion temperature is reflected in the profiles of the y -parameter shown in Figure 9. Here, the solid curves correspond to the location in the corona where $y = 1$, whereas the dashed curves are for $y = 0.5$. As one would expect, the corona appears more extended in those regions where $T_i \approx T_{vir}$. The temperature, T_e jumps to a value ≈ 50 keV near the boundary (at $R_0 \approx 15r_g$) where the instability first arises, and then continues to increase towards ≈ 350 keV as the matter

approaches $\sim 6 - 7r_g$. Thus, since most of the bremsstrahlung photons are emitted at $y > 1$, it is evident from Figure 5 that the location of the bump at ~ 1 MeV is a natural consequence of the self-Comptonization of the radiation, which approaches a Wien peak centered at roughly $3kT_e$ in this inner hot region. However, the midplane density ρ_0 is lower in the γ_2 and γ_3 states compared to γ_1 , and since the bremsstrahlung intensity scales as ρ_0^2 , the appearance of the bump is more evident in γ_1 than γ_2 and γ_3 .

Using these results, we may now justify *a posteriori* our neglect of pairs in the equilibrium lepton number density. Under the simplifying (and conservative) assumption that the radiation distribution within the corona is isotropic, we estimate the pair production rate via $\gamma\gamma$ interactions to be about $9 \times 10^{40} \text{ s}^{-1}$ in the γ_1 state. (The pair creation rate is smaller in the γ_2 and γ_3 states due to the relatively lower number density of $E > 511$ keV γ -rays.) The average (accretion) electron number density is $\langle n_e \rangle \approx 5 \times 10^{17} \text{ cm}^{-3}$, for which the positron (in-flight) annihilation time is $t_a \equiv 1/(n_e \sigma_T c) \approx 1 \times 10^{-4} \text{ s}$. This is in fact the shortest time scale associated with the positron distribution, providing an upper limit to the positron number density n_{pairs} . We therefore expect the positron average equilibrium number density to be $\langle n_{pairs} \rangle \lesssim 8 \times 10^{13} \text{ cm}^{-3}$ for the coronal structure depicted in Figure 5, in which the scale size is $\approx 3 \times 10^7 \text{ cm}$. As such, $n_{pairs} \ll n_e$ and our approach would appear to be valid, at least in the case of Cygnus X-1. Incidentally, if we further assume that the subsequent conversion of these pairs into 511 keV photons is an efficient one, the 511 keV luminosity should be $\approx 2.8 \times 10^{36} \text{ ergs s}^{-1}$, or about 10 times weaker than the luminosity in the bump.

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Figure Captions

Figure 1. Theoretical spectrum corresponding to the γ_1 state of Cygnus X-1. *Dashed curve:* bremsstrahlung/self-Compton component (BSC); *Solid curve:* overall spectrum, including the BSC component, the disk radiation upscattered within the $y \lesssim 1$ region of the inner corona, and the thermal disk radiation at $\lesssim 3$ keV. Data for $E \gtrsim 50$ keV are taken from Ling et al. 1987. Data below ~ 50 keV (which however were collected at a different epoc) are from Schreier et al. 1971.

Figure 2. Same as Figure 1, except now for the γ_2 state of Cygnus X-1. All data are from Ling et al. 1987.

Figure 3. Same as Figure 1, except now for the γ_3 state of Cygnus X-1. All data are from Ling et al. 1987.

Figure 4. Comparison of the γ_1 and γ_2 theoretical spectra, showing the cross-over at approximately 400 keV, and the anti-correlation between the strength of the bump at ≈ 1 MeV and the continuum flux below this cross-over energy.

Figure 5. Comparison of the γ_2 and γ_3 theoretical spectra, showing the overall shift in flux, for both the continuum and the bump, when \dot{M} changes by $\approx 25\%$.

Figure 6. The run of electron temperature T_e as a function of the radius r in the disk. The onset of the secular instability at $r \approx 15 - 17 r_g$ is signaled by a rapid increase in both T_e (to $\gtrsim 50$ keV), as specified by Equation (6) or (10). Because of gravitational stratification, the corona becomes optically thick toward smaller radii, and T_e increases further to enhance the cooling rate due to bremsstrahlung-self-Compton emission in this region.

Figure 7. Profiles of the mid-plane density as a function of radius for each of the three

gamma-ray states. The slight jump in density reflects a transition from the bound corona (for which $T_i < T_{vir}$) to a “coronal wind” zone (in which $T_i \approx T_{vir}$).

Figure 8. The run of ion temperature T_i as a function of radius for each of the three gamma-ray states. T_i is limited from above by the value of the virial temperature T_{vir} .

Figure 9. Contours of the Compton y -parameter for $y = 0.5$ (doted) and 1.0 (solid). (a) γ_1 state, (b) γ_2 state, and (c) γ_3 state. Distances are in units of the Schwarzschild radius $r_g \equiv 2GM/c^2$.